Where are the next Higgs bosons?

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Simple symmetry arguments applied to the third generation lead to a prediction: there exist new sequential Higgs doublets with masses of order $\lesssim 5.5$ TeV, with approximately universal Higgs-Yukawa coupling constants, $g \sim 1$. This is calibrated by the known Higgs boson mass, the top quark Higgs-Yukawa coupling, and the b -quark mass. A new massive weak-isodoublet, H_b , coupled to the b -quark with $g \sim 1$ is predicted. It may be accessible at the 13 TeV LHC, and definitively at an energy upgraded 26 TeV LHC. The extension to leptons generates a new H_τ and a possible H_{ν_τ} doublet. The accessibility of the latter depends upon whether the mass of the τ -neutrino is Dirac or Majorana. The new physics of the model is technically natural, and overall the model is no less natural than the Standard Model.

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I. INTRODUCTION

Understanding flavor physics will likely involve the discovery of new particles, associated with the mystery of the origin of the small parameters of the Standard Model. For example, the observed Higgs-Yukawa coupling of the b -quark is small, $y_b \simeq 0.024$. This is “technically natural” [1] because as $y_b \rightarrow 0$, a chiral symmetry emerges, $b_R \rightarrow e^{i\theta} b_R$, and a nonzero y_b is never perturbatively regenerated. However, the small parameters may have a perturbative origin, arising from virtual effects involving new particles with larger couplings; i.e., a small parameter starts as a large parameter that is subsequently power-law suppressed. Such new physics may be accessible to the LHC or its energy doubled upgrade.

There are, of course, many other theoretical ways to achieve this, but we are also motivated by the hypothesis that a single lone Higgs boson is unlikely to exist—there may be a rich spectrum of Higgs bosons, presenting a new spectroscopy in nature. Thus, the “new particles” we will consider are exclusively new massive Higgs isodoublets.

Hence, a plausible origin of the small y_b is via a new heavy Higgs isodoublet, H_b , coupled as $g_b \bar{T}_L H_b b_R$ where $T_L = (t, b)_L$. Most importantly, the new coupling g_b is large,

owing to a symmetry such as we propose below, where it is of order the top quark Higgs-Yukawa coupling, $g_b = O(1)$. The observed y_b is then “effective,” that is, $y_b \sim g_b(\mu^2/M_b^2)$, where the power-law suppression arises from the large mass term of the new H_b , $M_b^2 H_b^\dagger H_b$, and its mixing with the Standard Model Higgs doublet H_0 , $\mu^2 H_0^\dagger H_b + \text{H.c.}$, with $\mu^2/M_b^2 \ll 1$. The breaking of the b_R chiral symmetry is then governed, not by the small y_b , but rather by a large $g_b \sim 1$, yet the observed $y_b \sim 10^{-2}$ arises naturally. A UV completion with symmetries and nontrivial renormalization group evolution has a better chance of predicting the large g_b than the small y_b . One would hope to directly observe the heavy H_b at its mass, M_b . As a bonus, a larger g_b enhances the production and detection possibilities for H_b at the LHC or any other collider.

Presently we examine this scenario in detail, which we believe describes the first sequential new Higgs doublets that could emerge at the LHC. We will focus, for simplicity, upon the third generation, and we will ignore the masses and mixings with the first and second generations. However this approach, in an extended “scalar democracy,” can explain the remaining small Higgs-Yukawa couplings of the full Standard Model, while flavor constraints from Cabibbo-Kobayashi-Maskawa mixing and hierarchies are consistent [2].

We presently assume that the top-bottom subsystem is approximately invariant under a simple extension of the Standard Model symmetry group structure

$$G = \text{SU}(2)_L \times \text{SU}(2)_R \times \text{U}(1)_{B-L} \times \text{U}(1)_A. \quad (1)$$

By “approximately” we mean that, if we turn off the $U(1)_Y$ gauging, $g_1 \rightarrow 0$, the symmetry G is exact in the $d = 4$

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operators (kinetic terms, Higgs-Yukawa couplings, and potential terms). The Standard Model (SM) gauging is the usual $SU(2)_L \times U(1)_Y$ and is a subgroup of G . The $U(1)_Y$ generator is now $I_{3R} + (B - L)/2$. This electroweak gauging weakly breaks the symmetry $SU(2)_R \times U(1)_{B-L} \rightarrow U(1)_Y \times U(1)_{B-L}$; however, the $SU(2)_R$ remains as an approximate global symmetry of the $d = 4$ operators. In addition, a global $U(1)_A$ arises as well. G will provide custodial symmetry so that new symmetry breaking effects, in $d = 2$ operators, are technically natural. Beyond the usual naturalness issue of the Higgs boson mass, our present model is no less natural than the Standard Model.

To implement G in the (t, b) sector we require that the Standard Model Higgs doublet, H_0 , couple to t_R with coupling $y_t \equiv g_t$ in the usual way, and a second Higgs doublet, H_b , couple to b_R with coupling g_b . The symmetry G then dictates that there is only a single Higgs-Yukawa coupling $g = g_t = g_b$ in the quark sector. This coupling is thus determined by the known top quark Higgs-Yukawa coupling, $y_t = g \simeq 1$. A schematic proposal for a UV completion theory that leads to G , based upon far UV compositeness, was proposed in [2]. In that scheme $y_t = g \simeq g_b \simeq 1$ is actually predicted by the infrared fixed point [3,4]. Our current discussion is a simplified subsector of that larger theory. We do not presently require the ingredients of the larger theory, and we simply calibrate $g \simeq 1$ from experiment.

Since $m_b/m_t \ll 1$, the $SU(2)_R$ must be broken. Here we follow the old rules of chiral dynamics, deploying “soft” symmetry breaking through bosonic mass terms. The symmetry breaking in the $d = 2$ Higgs mass terms preserves the universality of quark Higgs-Yukawa couplings g . We remark that these explicit $d = 2$ symmetry-breaking mass terms could arise from $d = 4$ scale invariant interactions involving additional new fields and hence G could then be broken spontaneously. However, we are interested presently in the simplest phenomenological scheme and will be content to insert the $d = 2$ symmetry breaking mass terms by hand.

This simple (t, b) quark scenario with G is then predictive, because of the universal coupling. As we shall observe below, the natural mass scale for H_b is found to be $\lesssim 5.5$ TeV. This prediction involves the assumption that the input Higgs mass, M_H^2 , is very small, or by a “no-fine-tuning” argument similar to that of the Higgs mass in the minimal supersymmetric standard model.

Here we have three *a priori* unknown renormalized parameters, M_H^2 (the input Standard Model Higgs potential mass), μ^2 (the mixing of H_b with the Standard Model Higgs) and M_b^2 (the heavy H_b mass). We do not specify a UV completion and thus do not solve the naturalness problems of these quantities nor explain their origin. However, once these are specified, we obtain the b -quark mass $m_b \sim m_t(\mu^2/M_b^2)$ and the Higgs potential mass, $M_0^2 = -(88.4 \text{ GeV})^2 = M_H^2 - \mu^4/M_b^2$ (note that $|M_0| = 125 \text{ GeV}/\sqrt{2}$). The tachyonic

(negative) M_0^2 can thus be generated from an unknown input value, M_H^2 , but we should *not* fine-tune the difference, $M_H^2 - \mu^4/M_b^2$. This requirement establishes the scale of $M_b \sim 5$ TeV. Or, if we postulate that the input Higgs mass is small, $M_H^2 \ll M_b^2$, we immediately obtain $M_b \simeq 5.5$ TeV.

However, this result is modified when we extend the theory to include the third generation leptons (ν_τ, τ). This requires two additional Higgs fields, H_τ and H_{ν_τ} , with masses M_τ and M_ν . The presence of the third generation leptonic Higgs fields with an extension of the symmetry, G , has the general effect of reducing all the heavy Higgs masses, e.g., of H_b to ~ 3.6 TeV with H_τ , and possibly H_{ν_τ} , nearby.

Remarkably, the new Higgs spectrum is then systematically dependent upon whether neutrinos receive only Dirac masses, via a “Dirac-Higgs seesaw mechanism,” or Majorana masses in addition to Dirac masses through the type I seesaw mechanism. If neutrino masses are pure Dirac, the new H_{ν_τ} must be ultraheavy to produce a seesaw in μ^2/M_ν^2 , and only the H_τ is detectable; if on the other hand the neutrino mass is Majorana, then we expect $M_\nu \sim M_\tau$. Hence, the physics of neutrinos and the new Higgs spectrum are intimately interwoven here.

In view of the simplicity and natural symmetry basis of these results we will therefore focus on the third generation subsystem in some detail.

II. THE TOP-BOTTOM SUBSYSTEM

The assumption of the symmetry, G , of Eq. (1) for the top-bottom system leads to a Higgs-Yukawa (HY) structure that is reminiscent of the chiral Lagrangian of the proton and neutron (or the chiral constituent model of up and down quarks [5]; in a composite Higgs scenario based upon [6] this model was considered by Luty [7]):

$$V_{HY} = g \bar{\Psi}_L \Sigma \Psi_R + \text{H.c.} \quad \text{where } \Psi = \begin{pmatrix} t \\ b \end{pmatrix}. \quad (2)$$

In Eq. (2), Σ is a 2×2 complex matrix and V_{HY} is invariant under $\Sigma \rightarrow U_L \Sigma U_R^\dagger$ with $\Psi_L \rightarrow U_L \Psi_L$ and $\Psi_R \rightarrow U_R \Psi_R$. Note that the $U(1)_Y$ generator Y of the SM now becomes $Y = I_{3R} + (B - L)/2$ and furthermore Σ is neutral under $B - L$ which implies its weak hypercharge $[Y, \Sigma] = \Sigma I_{3R}$.

An additional $U(1)_A$ axial symmetry arises as an overall phase transformation of $\Sigma \rightarrow e^{i\theta} \Sigma$, accompanied by $\Psi_L \rightarrow e^{i\theta/2} \Psi_L$ and $\Psi_R \rightarrow e^{-i\theta/2} \Psi_R$. Keeping or breaking this symmetry is an arbitrary option for us [though in the (u, d) subsystem this is the Peccei-Quinn symmetry and is required if one incorporates the axion].

Σ can be written in terms of two column doublets,

$$\Sigma = (H_0, H_b^c), \quad (3)$$

where $H^c = i\sigma_2 H^*$. Under $SU(2)_L$ we have $H \rightarrow U_L H$ and $H^c \rightarrow U_L H^c$ where we note that the weak hypercharge

eigenvalues of H and H^c have opposite signs. The HY couplings of Eq. (2) become

$$V_{HY} = g_t(\bar{t}, \bar{b})_L H_0 t_R + g_b(\bar{t}, \bar{b})_L H_b^c b_R + \text{H.c.}, \quad (4)$$

where the $SU(2)_R$ symmetry has forced $g = g_t = g_b$.

Note that this becomes identical to the Standard Model if we make the identification

$$H_b \rightarrow \epsilon H_0, \quad \epsilon = \frac{m_b}{m_t} = 0.0234. \quad (5)$$

The HY coupling of the b -quark to H_0 is then $y_b = \epsilon g_b$, but with $\epsilon \ll 1$ the $SU(2)_R$ symmetry is then lost. The SM Higgs boson (SMH), H_0 , in the absence of H_b , has the usual SM potential:

$$V_{\text{Higgs}} = M_0^2 H_0^\dagger H_0 + \frac{\lambda}{2} (H_0^\dagger H_0)^2, \quad (6)$$

where $M_0^2 \simeq -(88.4) \text{ GeV}^2$ (note the sign), and $\lambda \simeq 0.25$. Minimizing this, we find that the Higgs field H_0 acquires its usual vacuum expectation value (VEV), $v = |M_0|/\sqrt{\lambda} = 174 \text{ GeV}$, and the observed physical Higgs boson, h , acquires mass $m_h = \sqrt{2}|M_0| \simeq 125 \text{ GeV}$. In our present scheme Eqs. (5) and (6) arise at low energies dynamically.

Given Eqs. (2) and (3), we postulate a new potential of the form

$$V = M_1^2 \text{Tr}(\Sigma^\dagger \Sigma) - M_2^2 \text{Tr}(\Sigma^\dagger \Sigma \sigma_3) + \mu^2 (e^{i\theta} \det \Sigma + \text{H.c.}) + \frac{\lambda_1}{2} \text{Tr}(\Sigma^\dagger \Sigma)^2 + \lambda_2 |\det \Sigma|^2. \quad (7)$$

We have written the potential in the Σ notation in order to display the symmetries more clearly. In the above, σ_3 acts on the $SU(2)_R$ side of Σ ; hence in the limit $M_2^2 = 0$ the potential V is invariant under $SU(2)_R$, while the μ^2 term breaks the additional $U(1)_A$. The associated CP -phase can be removed by field redefinition in the present model. If the exact $U(1)_A$ symmetry is imposed on the $d = 4$ terms in the potential then operators such as $e^{i\alpha}(\det \Sigma)\text{Tr}(\Sigma^\dagger \Sigma)$, $e^{i\alpha'}(\det \Sigma)^2$, etc., are forbidden. Noting the identity $(\text{Tr}(\Sigma^\dagger \Sigma))^2 = \text{Tr}(\Sigma^\dagger \Sigma)^2 + 2 \det \Sigma^\dagger \Sigma$, we obtain only the two indicated $d = 4$ terms as the maximal form of the invariant potential.

Hence the global symmetries of G restrict the $d = 4$ terms of this model. They also make the $d = 2$ symmetry breaking terms technically natural. For example, the $d = 4$ operator $\text{Tr}(\Sigma^\dagger \Sigma) \det \Sigma$ is disallowed by $U(1)_A$. However the $d = 2$ 't Hooft operator, $\mu^2 \det \Sigma$, breaks $U(1)_A$, is perturbatively multiplicatively renormalized and is technically naturally small (we are not considering nonperturbative instantons here).

The key point is that $d = 2$ operators do not renormalize $d = 4$ operators, and they may be effective operators that

parametrize new physics sectors. In fact, our theory is no less natural than the Standard Model. The analogue of the Higgs boson mass is the $d = 2$ operator, $\text{Tr}(\Sigma^\dagger \Sigma)$, which is allowed by G , but this is no more or less natural than the SM Higgs mass term. The other two $d = 2$ terms, $\mu^2 \det \Sigma$ and $\text{Tr}(\Sigma^\dagger \Sigma \sigma_z)$, explicitly break $U(1)_A$ and $SU(2)_R$, respectively; are multiplicatively renormalized; and can be technically naturally small.

Using Eq. (3), V can be written in terms of H_0 and H_b :

$$V = M_H^2 H_0^\dagger H_0 + M_b^2 H_b^\dagger H_b + \mu^2 (e^{i\theta} H_0^\dagger H_b + \text{H.c.}) + \frac{\lambda}{2} (H_0^\dagger H_0 + H_b^\dagger H_b)^2 + \lambda' (H_0^\dagger H_b H_b^\dagger H_0), \quad (8)$$

where $\lambda_2 = \lambda' + \lambda$, $M_H^2 = M_1^2 - M_2^2$, and $M_b^2 = M_1^2 + M_2^2$. In the limit $M_b^2 \rightarrow \infty$ the field H_b decouples, and Eq. (8) reduces to the SM if $M_H^2 \rightarrow M_0^2$, with $\lambda \rightarrow 0.25$ thereby recovering Eq. (6).

We assume the quartic couplings are all of order of the SM value $\lambda \sim 0.25$. Those associated with the new heavy Higgs bosons, such as $\lambda' \sim \lambda$, will therefore contribute negligibly small effects since they involve large, positive, M^2 heavy Higgs fields. We also set $\theta = 0$. Then, varying the potential with respect to H_b , the low momentum components of H_b become locked to H_0 :

$$H_b = -\frac{\mu^2}{M_b^2} H_0 + O(\lambda, \lambda'). \quad (9)$$

Substituting back into V we recover the SMH potential

$$V = M_0^2 H_0^\dagger H_0 + \frac{\lambda}{2} (H_0^\dagger H_0)^2 + O\left(\frac{\mu^2}{M_b^2}\right), \quad (10)$$

with

$$M_0^2 = M_H^2 - \frac{\mu^4}{M_b^2}. \quad (11)$$

Note that, even with M_H^2 positive, $M_0^2 = -(88.4)^2 \text{ GeV}^2$ can be driven to its negative value by the mixing with H_b (level repulsion). We minimize the SMH potential and define

$$H_0 = \begin{pmatrix} v + \frac{1}{\sqrt{2}} h \\ 0 \end{pmatrix}, \quad v = 174 \text{ GeV}, \quad (12)$$

in the unitary gauge. The minimum of Eq. (12) yields the usual SM result

$$v^2 = -M_0^2/\lambda, \quad m_h = \sqrt{2}|M_0| = 125 \text{ GeV}, \quad (13)$$

where m_h is the propagating Higgs boson mass. We can then write

$$H_b \rightarrow H_b - \frac{\mu^2}{M_b^2} \begin{pmatrix} v + \frac{1}{\sqrt{2}}h \\ 0 \end{pmatrix}, \quad (14)$$

for the full H_b field. This is a linearized (small angle) approximation to the mixing and is reasonably insensitive to the small $\lambda, \lambda' \ll 1$.

Note the effect of “level repulsion” of the Higgs mass, M_0^2 , downward due to the mixing with heavier H_b . The level repulsion in the presence of μ^2 and M_b^2 occurs due to an approximate “seesaw” Higgs mass matrix

$$\begin{pmatrix} M_H^2 & \mu^2 \\ \mu^2 & M_b^2 \end{pmatrix}. \quad (15)$$

The input value of the mass term M_H^2 is unknown and in principle arbitrary, and it can have either sign. We can presumably bound $M_H^2 \gtrsim 1 \text{ GeV}^2$ from below, since QCD effects will mix glueballs and the QCD- σ -meson with H_0 in this limit. As M_H^2 is otherwise arbitrary, we might then expect that the most probable value is $M_H^2 \sim 1 \text{ GeV}^2$, and $|M_H^2| \ll (\mu^2, M_b^2)$.

Let us consider the case $M_b^2 \gg M_H^2$. Then Eq. (15) has eigenvalues $M_0^2 = -\mu^4/M_b^2$ and M_b^2 . Thus, in the limit of small, nonzero $|M_H^2|$, we see that a negative M_0^2 arises naturally, and to a good approximation the physical Higgs mass is generated entirely by this negative mixing term.

The mass mixing causes the neutral component of H_b to acquire a small VEV (“tadpole”) of $-v(\mu^2/M_b^2)$. This implies that the SM HY-coupling of the b -quark is induced with the small value $y_b = g_b(\mu^2/M_b^2)$ (note that g_b is positive with a phase redefinition of b_R). The b -quark then receives its mass from H_b ,

$$m_b = g_b(m_b)v \frac{\mu^2}{M_b^2} = m_t \frac{g_b(m_b)\mu^2}{g_t(m_t)M_b^2}, \quad (16)$$

where we have indicated the renormalization group (RG) scales at which these couplings should be evaluated.

In a larger framework with a UV completion, such as Ref. [2], at a mass scale, $m > M_b$, both $g_t(m)$ and $g_b(m)$ will have a common renormalization group equation modulo $U(1)_Y$ effects. We implicitly assume that, at some very high scale $\Lambda \gg v$, the $SU(2)_L \times SU(2)_R$ is a good symmetry. This implies $g(\Lambda) = g_t(\Lambda) = g_b(\Lambda) \gtrsim 1$. Then, we will predict $g(M_b) \simeq g_t(M_b) \simeq g_b(M_b)$, e.g., where the values at the mass scale M_b are determined by the RG fixed point [3,4] with small splittings due to $U(1)_Y$.

Furthermore, we find that $g_t(m)$ and, more so, $g_b(m)$ increase somewhat as we evolve downward from M_b to m_t or m_b . The top quark mass is then $m_t = g_t(m_t)v$ where v is the SM Higgs VEV. From these effects we obtain the ratio

$$R_b = \frac{g_b(m_b)}{g_t(m_t)} \simeq 1.5. \quad (17)$$

The b -quark then receives its mass from the tadpole VEV of H_b :

$$m_b = g_b(m_b)v \frac{\mu^2}{M_b^2} = m_t R_b \frac{\mu^2}{M_b^2}. \quad (18)$$

In the case that the Higgs mass, M_0^2 , is due entirely to the level repulsion by H_b , i.e., $M_H^2 = 0$, and using Eqs. (11), (17) and (18), we obtain a predicted mass of the H_b ,

$$M_b = \frac{m_t}{m_b} R_b |M_0| \simeq 5.5 \text{ TeV}, \quad (19)$$

with $m_b = 4.18 \text{ GeV}$, $m_t = 173 \text{ GeV}$, and $|M_0| = 88.4 \text{ GeV}$. We remind the reader that we have ignored the effects of the quartic couplings λ , which we expect are small. Moreover, the quartic couplings do not enter the mixing, because terms such as $H_0^\dagger H_0 H_0^\dagger H_b$ are forbidden by our symmetry. The remaining terms only act as slight shifts in the masses, never larger than $\sim \lambda v^2$, and can be safely ignored.

This is a key prediction of the model (see also [8,9]). In fact, we can argue that with M_H^2 nonzero, but with small fine-tuning (see below), the result $M_b \lesssim 5.5 \text{ TeV}$ is obtained. This mass scale is accessible to the LHC with luminosity and energy upgrades, and we feel that it represents an important target for discovery of the first sequential Higgs boson.

However, we will see in the next section that this result is reduced when the third generation leptons are included.

III. INCLUSION OF THE (ν_τ, τ) LEPTONS

The simple (t, b) system described above can be extended to the third generation leptons (ν_τ, τ) . Presently we will abbreviate ν_τ to ν and ignore neutrino mixing in this third generation scheme (see also [10]). We consider two distinct mechanisms to generate the small neutrino mass scale.

Remarkably the predictions for the mass spectrum are sensitive to the mechanism of neutrino mass generation. The next sequential massive Higgs isodoublet, in addition to H_b , is likely to include the H_τ , and possibly also H_ν , which is dependent upon whether neutrino masses are Majorana or Dirac in nature.

We introduce the Higgs-Yukawa couplings for the leptons in a G invariant form

$$\begin{aligned} V_{HY\ell} &= g_\ell \bar{\Psi}_L \Sigma_\ell \Psi_R + \text{H.c.} \\ &= g_\ell (\bar{\nu}, \bar{\tau})_L H_\nu \nu_R + g_\tau (\bar{\nu}, \bar{\tau})_L H_\tau^c \tau_R + \text{H.c.}, \end{aligned} \quad (20)$$

where we have introduced a second “leptonic” chiral field

$$\Sigma_\ell = (H_\nu, H_\tau^c), \quad \text{where } \Psi = \begin{pmatrix} \nu \\ \tau \end{pmatrix}, \quad (21)$$

and Σ_ℓ transforms as Σ under G .

The couplings g_ν and g_τ are assumed to have the common value at the high scale Λ ,

$$g = g_t(\Lambda) = g_b(\Lambda) = g_\tau(\Lambda) = g_\nu(\Lambda). \quad (22)$$

Since the RG equations for the leptons do not involve QCD, we typically find $g_\tau(m_\tau) = g_\nu(m_\tau) \simeq 0.7$ for $\Lambda \sim M_{\text{Planck}}$ [2]. More generally we define the parameters

$$R_\tau = \frac{g_\tau(m_\tau)}{g_t(m_t)} \quad R_\nu = \frac{g_\tau(m_\nu)}{g_t(m_t)}. \quad (23)$$

We extend the potential of Eq. (7), $V \rightarrow V + V'$ to incorporate the Σ_ℓ mass terms:

$$\begin{aligned} V'_0 &= M_\ell^2 \text{Tr}(\Sigma_\ell^\dagger \Sigma_\ell) + \delta M_\ell^2 \text{Tr}(\Sigma_\ell^\dagger \Sigma_\ell \sigma_z) \\ &= M_\nu^2 H_\nu^\dagger H_\nu + M_\tau^2 H_\tau^\dagger H_\tau + \dots \end{aligned} \quad (24)$$

where the ellipsis refers to quartic terms which we will ignore altogether. Moreover, we assume that the new isodoublets are dormant, $M_\nu^2, M_\tau^2 > 0$.

We can also introduce a mixed term that involves both Σ_ℓ and Σ of Eq. (3):

$$\begin{aligned} V'_1 &= \mu_1^2 \text{Tr}(e^{i\theta'} \Sigma^\dagger \Sigma_\ell + \text{H.c.}) \\ &= \mu_1^2 e^{i\theta'} (H_\nu^\dagger H_0 + H_b^\dagger H_\tau) + \text{H.c.} \end{aligned} \quad (25)$$

Note that this term mixes the Higgs fields as $H_0 \leftrightarrow H_\nu$ and $H_b \leftrightarrow H_\tau$. Such mixing would lead to the τ acquiring its mass sequentially via mixing with H_b , which directly mixes with H_0 as in Eq. (8). Although such a scenario has potentially interesting physics, we do not pursue it presently.

However, we can introduce a second term consistent with G that leads to direct mixing of $H_0^\dagger H_\tau$. This can be constructed using a charge conjugated Σ^c , where

$$\Sigma_\ell^c = i\sigma_2 \Sigma_\ell^* (-i\sigma_2), \quad (26)$$

and note that $\Sigma^c \rightarrow U_L \Sigma^c U_R$ transforms identically to $\Sigma \rightarrow U_L \Sigma U_R$ under the $SU(2)$ groups. The effect of the conjugation is to flip the column and charge conjugation assignments of Higgs isodoublets in Σ_ℓ ,

$$\Sigma_\ell^c = (H_\tau, H_\nu^c). \quad (27)$$

Therefore, a term that permits direct mixing $H_0 \leftrightarrow H_\tau$ and $H_b \leftrightarrow H_\nu$ is

$$\begin{aligned} V'_2 &= \mu_2^2 e^{i\phi'} \text{Tr}(\Sigma_\ell^c \Sigma) + \text{H.c.} \\ &= \mu_2^2 e^{i\phi'} (H_\tau^\dagger H_0 + H_b^\dagger H_\nu) + \text{H.c.} \end{aligned} \quad (28)$$

It should be noted that this term violates the $U(1)_A$ symmetry, but trivially not $U(1)_{B-L}$ since both Σ 's are sterile under $B-L$. For simplicity we will simply set the CP phases to zero, $\theta = \theta' = \phi' = 0$.

A. H_τ , H_b , and a Dirac neutrino seesaw

One interesting possibility is that the $SU(2)_R$ symmetry is badly broken in the lepton sector with $M_\nu^2 \gg M_\tau^2$ of Eq. (24). In this limit the H_ν has become ultramassive and nondetectable at collider energies.

Through the V'_1 term we observe that H_ν acquires a tiny tadpole VEV, but now has negligible feedback on the Higgs mass:

$$H_\nu = -\frac{\mu_1^2}{M_\nu^2} H_0, \quad \delta M_0^2 = -\frac{\mu_1^4}{M_\nu^2} \simeq 0, \quad (29)$$

where $\mu_1^2/M_\nu^2 \ll 1$. Hence the neutrino acquires a tiny Dirac mass through the induced coupling to the Higgs:

$$\begin{aligned} V_{HY\ell} &= -g_\nu \frac{\mu_1^2}{M_\nu^2} \left(\bar{\nu} \right)_L H_0 \nu_R + \dots \\ m_\nu &= g_\nu v \frac{\mu_1^2}{M_\nu^2} = m_t R_\nu \frac{\mu_1^2}{M_\nu^2}, \end{aligned} \quad (30)$$

where $R_\nu \simeq O(1)$. In such a scenario a Majorana mass term is not necessary, as we have a (Dirac) seesaw mechanism to naturally generate a tiny neutrino mass.

The H_τ mixes directly to the SMH through the V'_2 term. H_τ then acquires a VEV and feeds back upon the Higgs mass as

$$H_\tau = -\frac{\mu_2^2}{M_\tau^2} H_0, \quad \delta M_0^2 = -\frac{\mu_2^4}{M_\tau^2}. \quad (31)$$

Hence, the τ mass is given by

$$m_\tau = m_t R_\tau \frac{\mu_2^2}{M_\tau^2} \quad (32)$$

and we expect $R_\tau \sim 0.7$.

H_τ and H_b now simultaneously contribute to the SMH mass

$$M_0^2 = M_H^2 - \frac{\mu^4}{M_b^2} - \frac{\mu_2^4}{M_\tau^2}. \quad (33)$$

[We remind the reader that $M_0^2 = -(88.4)^2 \text{ GeV}^2$ is negative as defined in Eq. (6)]. Using Eqs. (16) and (32) this yields an elliptical constraint on the heavy Higgs masses M_b and M_τ ,

FIG. 1. The fine-tuning associated with different values of the M_b and M_τ parameters. The remaining three parameters μ , μ_τ , and M_H are fixed by the physical choices of fermion masses and Higgs VEV. The red dashed line corresponds to $M_H^2 = 0$ and the origin to $M_H^2 = -|M_0^2|$.

$$|M_0|^2 + M_H^2 = \frac{m_b^2 M_b^2}{m_\tau^2 R_b^2} + \frac{m_\tau^2 M_\tau^2}{m_\tau^2 R_\tau^2}, \quad (34)$$

for fixed M_H^2 . Bear in mind that the Higgs potential input mass, M_H^2 , is *a priori* unknown, while $M_0^2 = -(88.4 \text{ GeV})^2$ is known from the Higgs boson mass, $m_h = \sqrt{2}|M_0| = 125 \text{ GeV}$.

If we make the assumption $M_H^2 = 0$ the ellipse is shown as the red dashed line in Fig. 1. If we further assume that both H_b and H_τ contribute equally to the SMH mass, then we obtain from Eq. (34)

$$M_b \simeq 3.6 \text{ TeV}, \quad M_\tau \simeq 4.2 \text{ TeV}. \quad (35)$$

Of course, we can raise (lower) these masses by introducing the bare positive (negative) M_H^2 . However, we do not want to excessively fine-tune the difference $-|M_0^2| + M_H^2$.

As an alternative way of estimating the Higgs masses we follow the same procedure for estimating the fine-tuning as is sometimes used in the minimal supersymmetric standard model or composite Higgs models. To account for the contribution to the fine-tuning from all parameters x_i on an observable O we use the measure (see e.g., [11])

$$\Delta = \left| \frac{\partial \ln O}{\partial \ln x_i} \right|. \quad (36)$$

That is to say, the fine-tuning is the length of the gradient of the logarithmic observable in the space of logarithmic parameters (other fine-tuning definitions have been explored and lead to quantitatively similar results).

Considering the case of two dormant Higgs bosons, H_b and H_τ , we measure the log-derivative sensitivities of the known value of M_0^2 to the five input parameters M_H^2 , μ_b^2 , μ_τ^2 , M_b^2 and M_τ^2 . The quadrature combined sensitivities yield the fine-tuning parameter, Δ . The fine-tuning is summarized in Fig. 1 where the red dashed curve indicates $M_H^2 = 0$. Every point on this figure is a bona fide fit to the known values of M_0^2 , m_b and m_τ .

Examining just the theory with H_b alone, we can combine the log-derivative sensitivities of M_0^2 , m_b to three input parameters M_H^2 , μ_b^2 , and M_b^2 . Then we find that the “no-fine-tuning limit” $\Delta \leq 1$ translates into the more restrictive constraint

$$M_b \leq 3.1 \text{ TeV}. \quad (37)$$

The lightest values of M_b (and M_τ) correspond to $-|M_0^2| = M_H^2$ with no contribution from M_b^2 (M_τ^2). Indeed, with more heavy Higgs bosons we have generically smaller masses and the various new particles become more easily accessible to the LHC.

B. H_ν , H_τ , H_b , and Majorana neutrinos

We can also consider the more conventional possibility that the neutrino has a Majorana mass term, which is an extension of the physics beyond the minimal model. The lepton sector Higgs masses may then be comparable, $M_\nu^2 \sim M_\tau^2$, and we have direct coupling to the SMH by both H_ν and H_τ , through the μ_1^2 and μ_2^2 terms. In the absence of a Majorana mass term, this would imply that G is a good symmetry, with the neutrino having a large Dirac mass $m_\nu = m_{D\nu}$ comparable to m_τ .

However, we can then suppress the physical neutrino mass, m_ν , by allowing the usual type I seesaw. We thus postulate a large Majorana mass term for the ungauged ν_R :

$$M \bar{\nu}_R^c \nu_R + \text{H.c.} \quad (38)$$

Integrating out ν_R we then have an induced $d = 5$ operator that generates a Majorana mass term for the left-handed neutrinos [12] through the VEV of H_ν ,

$$V_M = \frac{g^2}{2M} \bar{\Psi}_L \Sigma_\ell (\mathbb{1} - \sigma_3) \Sigma_\ell^c \Psi_L^c + \text{H.c.} \quad (39)$$

In this scenario there is no restriction that requires the H_ν mass to be heavy, and the neutrino physical mass is now small, given by $m_\nu \sim m_{D\nu}^2/M \sim m_\tau^2/M$ where $m_{D\nu} \sim m_\tau$ is the neutrino Dirac mass.

Through the V_1' term we find that H_ν acquires a VEV and has comparable feedback on the Higgs mass as H_τ ,

$$H_\nu = -\frac{\mu_1^2}{M_\nu^2} H_0, \quad \delta M_0^2 = -\frac{\mu_1^4}{M_\nu^2}. \quad (40)$$

We now have an ellipsoid for the masses M_b , M_τ and M_ν :

$$|M_0|^2 + M_H^2 = \frac{m_b^2 M_b^2}{m_\tau^2 R_b^2} + \frac{m_\tau^2 M_\tau^2}{m_t^2 R_\tau^2} + \frac{m_{D\nu}^2 M_\nu^2}{m_t^2 R_\nu^2}. \quad (41)$$

With additional light Higgs fields, the elliptical constraint forces all of the Higgs masses to smaller values. We emphasize that the mass bounds of Fig. 1 should be viewed as upper limits on the Higgs mass spectrum that could be explored at the LHC.

IV. PHENOMENOLOGY OF THE SEQUENTIAL HIGGS BOSONS

Presently we touch upon the collider phenomenology of this model and refer the reader to [2] for further discussion. Flavor physics bounds have also been considered in [2]. Our estimated discovery luminosities for h_b^0 are seen to be attainable at the LHC or its upgrades, and we thus encourage and plan more detailed studies. Moreover, the lower mass range $\lesssim 1$ TeV is currently within range of the LHC and collaborations should attempt to place limits.

H_b is an isodoublet with neutral h_b^0 and charged $h_b^\pm$ complex field components coupling to (t, b) as $g_b h_b^0 \bar{b}_L b_R + \text{H.c.}$ and $g_b h_b^\pm \bar{t}_L b_R + \text{H.c.}$ with $g_b \simeq 1$. At the LHC the h_b^0 is singly produced in pp by the perturbative (“intrinsic”) b and $\bar{b}$ components of the proton,

$$pp(b\bar{b}) \rightarrow h_b^0 \rightarrow b\bar{b}. \quad (42)$$

Structure functions that contain the $b + \bar{b}$ components are available in MadGraph5_aMC@NLO [13] and CalCHEP [14], and the results obtained by these two simulators are found to be consistent. The resulting cross sections for 13, 26, and 100 TeV are given for various M_b in [2] and in Fig. 2.

The cross section for production at the LHC of h_b^0 is $\sigma(h_b^0 \rightarrow b\bar{b}) \sim 10^{-4}$ pb at 13 TeV [or $\sigma(h_b^0 \rightarrow b\bar{b}) \sim 10^{-2}$ pb at 26 TeV] for a mass of $M_b = 3.5$ TeV. The decay width of h_b^0 is large, $\Gamma = 3M_b/16\pi \simeq 210$ GeV, for $M_b = 3.5$ TeV. To reduce the backgrounds we impose a 100 GeV p_T cut on the each b jet in the quoted cross sections. Note that the charged $h_b^\pm$ would be produced in association with $\bar{t}b$, with a significantly smaller cross section, and we have not analyzed it.

We note that the h^0 is centrally produced in a narrow range of rapidity, $|\eta| < 1$, which may afford useful cuts, though this is somewhat redundant to the 100 GeV p_T we used in this study.

The main backgrounds are high mass b -quark-dijet production $pp \rightarrow \bar{b}b$ and ordinary flavor dijets $pp \rightarrow \bar{q}q$ that fake $\bar{b}b$. The $pp \rightarrow \bar{b}b$ is mostly forward and requires the p_T cut, chosen to be 100 GeV/ c . We have not extensively studied the optimization of this choice for this cut. We use the same simulators to generate the backgrounds.

The light quark and gluon jets faking b -dijets can be viewed as a multiplier of this background. We find that as

FIG. 2. The left plot shows the estimated cross sections for $pp \rightarrow X + H_b(\rightarrow \bar{b}b)$ at the LHC 13 TeV, LHC upgraded 26 TeV, and a 100 TeV pp collider. The cross sections were calculated using structure functions with perturbative b and $\bar{b}$, applying a 100 GeV p_T cut on the final state b -jets, using MadGraph5_aMC@NLO [13] and CalCHEP [14]. The right plot displays the estimated 5σ integrated luminosities for H_b in the $\bar{b}b$ final state, applying a Breit-Wigner as described in the text, and assuming a 50% double b -tagging efficiency.

the b -dijet fake rate becomes smaller than $\epsilon \sim 3\%$ this becomes negligible. Single b -jet fake rates approaching $\epsilon \sim 1\%$ have been achieved in Tevatron DZero data [15], so suppressing fakes of b -dijets to a level of $\epsilon^2 < 10\%$ should be straightforward. We also note that there are other possible optimizations, such as the selection of the Breit-Wigner width (signal bin width) that we took to be $2\Gamma_{H_b}$. A central rapidity cut was not applied since it is somewhat redundant to the high mass and p_T cut.

We estimate that a 5σ excess in $S/\sqrt{B}$ for $M_b = 3.5$ TeV, in bins spanning twice the full width of the Breit-Wigner, requires an integrated luminosity of $\sim 20 \text{ ab}^{-1}$ at 13 TeV, or $\sim 100 \text{ fb}^{-1}$ at 26 TeV.¹ Similarly, at a mass of 5 TeV one requires $\sim 3 \text{ ab}^{-1}$ at 26 TeV for a 5σ discovery. This assumes double b -tagging efficiencies of order 50%. Here the background is assumed to be mainly $gg \rightarrow \bar{b}b$, but the large fake rate from $gg \rightarrow \bar{q}q$ must be reduced to $\lesssim 3\%$ for this to apply.

Hence according to our estimates, while meaningful bounds are achievable at the current 13 TeV LHC, particularly $M_b \lesssim 3.5$ TeV, the energy doubler is certainly favored for this physics and could cover the full mass range up to ~ 7 TeV. Note that we have not attempted any significant optimization of the $S/\sqrt{B}$ in this study, and it may require a detector dependent analysis to do significantly better.

Remarkably, the h_τ^0 (and h_ν^0), neutral components of the associated isodoublets, H_τ and H_ν , may also be singly produced because they can mix with H_b through the μ_1^2 and μ_2^2 terms of Eqs. (25) and (28) respectively. This implies a total cross section

$$\sigma(h_\tau^0, h_\nu^0) \sim \theta^2 \sigma(h_b^0), \quad (43)$$

for the mixing angle, θ , between either the states h_τ^0 or h_ν^0 and h_b^0 . Although θ is unknown it could easily be large, $\theta \sim 0.3$. The $h_\tau^0 \rightarrow \bar{\tau}\tau$ is visible with τ -tagging, and the background is also slightly suppressed since the peak is narrower by a factor of $(g_\tau^2/3g_b^2) \sim 0.16$. Hence at the 26 TeV LHC discovery is in principle possible for a 3.5 TeV state with an integrated luminosity of order 2 ab^{-1} .

These states are also pair produced by electroweak processes at the LHC or through $\gamma + Z^0$ at a high energy lepton collider:

$$\begin{aligned} \ell^+ \ell^- &\rightarrow (\gamma + Z^0)^* \rightarrow H_\tau^0 H_\tau^0 \rightarrow 4\tau \\ \ell^+ \ell^- &\rightarrow (\gamma + Z^0)^* \rightarrow H_x^\pm H_x^\mp \rightarrow 2\tau + \text{missing } p_T \\ pp &\rightarrow (W^\pm)^* \rightarrow H_\tau^\pm H_\tau^0 \rightarrow 3\tau + \text{missing } p_T \\ pp &\rightarrow (W^\pm)^* \rightarrow H_\nu^\pm H_\nu^0 \rightarrow \tau + \text{missing } p_T \end{aligned} \quad (44)$$

where x denotes either τ or ν .

We note that the LHC now has the capability of ruling out an H_b with $g_b \sim 1$ of mass ~ 1 TeV, with current integrated luminosities, $\sim 200 \text{ fb}^{-1}$.

V. CONCLUSIONS

We believe it is likely that an extended spectrum of Higgs bosons exists in nature. The expected masses of isodoublet Higgs bosons have been discussed here with the prominent new field, H_b , coupled to b_R in the range $\lesssim 5.5$ TeV. We emphasize that this is an approximate upper limit, and experiments should search and quote limits in the full mass range for H_b . Additional new states, such as leptonic Higgs isodoublets, or an extended spectroscopy as in [2], may emerge.

We strongly emphasize that the LHC *already has* the capability of ruling out an H_b with $g_b \sim 1$ of mass ~ 1 TeV, with current integrated luminosities, $\sim 200 \text{ fb}^{-1}$. The main reason is that the very plausible $g_b \simeq 1$ significantly enhances access to these states. We are unaware of any limits for these in the literature to date. We think it is important for the collaborations to develop an analysis strategy for limiting or discovery of these states.

This multi-Higgs scenario connects with the traditional way in which physics has evolved from atoms to nuclei to hadrons, thus presenting yet another spectroscopy [16]. Observation of the H_b , and measurement of $g_b \simeq 1$, would constitute compelling evidence of such an extended Higgs spectrum and the validity of the logic behind these schemes.

Here we have shown that a subsector of a more general scalar democracy [2] can be described simply in the context of the restricted, yet most observable, third generation by applying a simple extension of the Standard Model symmetry group, G . We think the overall simplicity of this idea is compelling, and we emphasize that the keystone is the universal Higgs-Yukawa coupling of order unity, controlled by a symmetry.

Above all, this theory is testable and should provide motivation to go further and deeper into the energy frontier with LHC upgrades, possibly with a future machine at the ~ 100 TeV scale and/or high energy lepton colliders. This scenario also suggests remarkable synergies with the ultraweak scale of neutrinos. For example, if the H_ν were inferred to exist with a mass in the $\lesssim 10$ TeV range, then neutrinos must have Majorana masses.

We note that the idea that small effective Higgs-Yukawa couplings are dynamical, as $y \sim g(\mu^2/M^2)$, is an old idea inherited from, e.g., extended technicolor models [17,18]. The primacy of the top quark in this scheme, with the calibrating large Higgs-Yukawa coupling, $g \simeq 1$, was anticipated long ago from considerations of the top infrared quasifixed point [3,4]. The fixed point corresponds to a Landau pole in y_t [4] and the Higgs then naturally arises as a $\bar{t}t$ composite state [6,19,20] (i.e., the fixed point [4] is the solution to top condensation models). The model of [2],

¹1 ab = 1 attobarn = 10^{-3} fb (femtobarn).

which contains the present model as a subgroup, has the SMH as a composite $H^0 \sim \bar{T}t$, and $H_b \sim \bar{T}b$, $H_\tau \sim \bar{L}\tau$, etc. The fixed point prediction becomes concordant with the top quark mass in our extended scheme and even probes the spectrum of Higgs bosons.

Many of the structural features of our potentials are similar to a chiral constituent quark model [5], though presently Σ is subcritical, with only the SM Higgs condensing. Indeed, the μ^2 mixing term above that generates the b -quark mass is equivalent to a 't Hooft determinant in chiral quark models, or in topcolor models [21].

Extended Higgs models are numerous, but those that most closely presage our present discussion are [8–10, 21–23]. In [2] we suggested a universal composite system of scalars, which we dubbed “scalar democracy.” However, in the more minimal third generation scenario presented here, which is most important for experimental observation, these features are determined only by the symmetry group G . This mainly makes all of the HY couplings of quarks and leptons universal, modulo RG running, and “hyperfine splitting” by the small SM gauge interactions. We also find it compelling that a negative M_0^2 mass can be generated by this mixing effect starting from a positive or very small input M_H^2 .

The universal value of all HY couplings, $g \simeq 1$, is of central importance to these models, and also enhances their

predictivity. The observed small parameters of the SM, such as y_b , y_τ , etc., are given essentially by one large universal parameter g multiplied by a suppressing power law, $\sim \mu^2/M^2$, together with smaller perturbative renormalization effects R_χ . In this way, the technically natural small couplings, and in principle the full Cabibbo-Kobayashi-Maskawa structure [2], can be understood to emerge from the UV and the inverted spectrum of Higgs bosons, and they may possibly be accessible to experiment.

A large extended spectrum of Higgs bosons with symmetry properties that unify all HY couplings to large common values is a phenomenologically rich idea, worthy of further theoretical study and the development of search strategies at the LHC and other future high energy colliders.

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